Gaugino condensation and the anomalous $U(1)$

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Abstract

We study gaugino condensation in presence of an anomalous $U(1)$ gauge group and find that global supersymmetry is dynamically broken. An example of particular interest is provided by effective string models with 4-dimensional Green-Schwarz anomaly cancellation mechanism. The structure of the hidden sector is constrained by the anomaly cancellation conditions and the scale of gaugino condensation is shifted compared with the usual case. We explicitly compute the resulting soft supersymmetry breaking terms.

LPTHE-Orsay 96/60

Saclay T96/075

hep-ph/9607172

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1 Introduction

Among the scenarios for breaking supersymmetry, the condensation of gauginos in a hidden sector remains a favourite one [1]. Gaugino condensation is indeed central to the idea of dynamical supersymmetry breaking in the context of gauge symmetries. But in its explicit realizations in superstring models [2], it suffers from a number of drawbacks. The degeneracies associated with the flat directions of the scalar potential are lifted but the corresponding degrees of freedom such as the dilaton are not stabilized and the true ground state is found at infinite field values where supersymmetry is restored. In order to overcome this problem, so-called "racetrack" or multicondensate models [3] have been proposed where two terms of different orders conspire in order to stabilise the field and provide a supersymmetry-breaking minimum.

Many of the superstring models have an anomalous $U(1)_X$ gauge symmetry [4] which could play a important role in issues such as fermion mass hierarchies [5, 6], cosmology [7] and ... gaugino condensation. The latter connection was recently stressed by Banks and Dine [8].[2] Indeed, the anomalous $U(1)_X$ has mixed anomalies with the other gauge symmetries –those of the standard model as well as the hidden sector–, anomalies which are cancelled through a 4-dimensional Green-Schwarz mechanism [9] using the couplings of the dilaton superfield to the gauge superfields. It is therefore not surprising that the whole issue of supersymmetry breaking through gaugino condensation is deeply modifed in such models. Moreover, because in the Green-Schwarz mechanism all the mixed anomalies are non-vanishing and proportional to one another, there must exist fields charged under $U(1)_X$ in the observable as well as in the hidden sector. The $U(1)_X$ gauge symmetry thus serves as a messenger interaction competitive with the gravitational interaction. In this paper, we wish to stress the modifications that the presence of such an anomalous $U(1)_X$ symmetry is bringing to the scenario of a

[2] We wish to thank Luis Ibáñez for drawing our attention to this paper and the whole issue.

dynamical supersymmetry breaking through gaugino condensation.

To be more specific, the relevant couplings of the dilaton superfield S to the gauge superfields V_a, V_X (of respective gauge invariant field strengths W_a^α, W_X^α) of the groups G_a and $U(1)_X$ read in the global limit:

$$\begin{aligned}\mathcal{L}_{S,V} &= -\int d^4\theta \ln(S+S^+ - \delta_{GS}V_X) \\ &\quad + \int d^2\theta \left[\frac{S}{4}(\sum_a k_a \mathrm{Tr}W_a^\alpha W_{a\alpha} + k_X \mathrm{Tr}W_X^\alpha W_{X\alpha}) + \mathrm{h.c.}\right] \end{aligned} \tag{1.1}$$

where δ_{GS} is the Green-Schwarz coefficient and k_a (k_X) is the Kac-Moody level of the group G_a $(U(1)_X)$. Under a $U(1)_X$ gauge transformation $(A_\mu^X \to A_\mu^X + \partial_\mu\alpha)$, S is shifted as

$$S \to S + \frac{i}{2}\delta_{GS}\alpha(x). \tag{1.2}$$

The complete Lagrangian is invariant provided the mixed $U(1)_X[G_a]^2$ anomaly coefficients C_a satisfy the condition

$$\delta_{GS} = \frac{C_a}{k_a} = \frac{C_X}{k_X} = \frac{C_g}{k_g}, \tag{1.3}$$

where C_g is the mixed gravitational anomaly proportional to $\mathrm{Tr}X$. Indeed a string computation yields

$$\delta_{GS} = \frac{1}{192\pi^2}\mathrm{Tr}X. \tag{1.4}$$

The mixing between S and V_X in the Kähler potential (1.1) gives rise to a D-term in the scalar potential:

$$V_D = \frac{g_X^2}{2}\left(\sum_A X_A K_A \phi^A + \frac{1}{4}k_X g_X^2 \delta_{GS} M_P^2\right)^2, \tag{1.5}$$

where M_P is the Planck scale,

$$k_X g_X^2 = \frac{2}{S+S^+}, \tag{1.6}$$

and K_A is the derivative of the Kähler potential K with respect to the field ϕ^A. The presence of the Fayet-Iliopoulos term induced by the $U(1)_X$

anomaly usually induces a non-zero vacuum expectation value for one (or more) field of X charge of sign opposite to δ_{GS}.

In the case where there is no anomalous $U(1)$, a non-anomalous R-symmetry under which the S superfield undergoes a translation similar to (1.2) imposes that the scale Λ at which the hidden sector gauge coupling becomes strong (which sets the scale for the corresponding gaugino condensates) behaves as

$$\Lambda \sim M_P e^{-\frac{kS}{2b_0}}, \tag{1.7}$$

where b_0 is the one-loop beta function coefficient for the hidden sector gauge group and k is its Kac-Moody level. This is obviously not invariant under the $U(1)_X$ transformation (1.2) of S. This shows that the presence of an anomalous $U(1)_X$ necessarily modifies the standard discussion of gaugino condensation.

This study goes beyond the case of effective superstring models. Indeed, it is not unusual that, in the course of sequential gauge symmetry breaking, appears an anomalous $U(1)_X$ abelian gauge symmetry whose mixed anomalies are cancelled through a Green-Schwarz type mechanism using an effective degree of freedom with dilaton-axion couplings. Thus we will consider in what follows the general case of a gauge model with symmetry group $SU(N_c) \times U(1)_X$ with a non-vanishing Fayet-Iliopoulos term and the following matter content: $N_f \leq N_c$ flavors, a dilaton-axion superfield and a chiral supermultiplet which breaks the anomalous $U(1)_X$ symmetry and whose vacuum expectation value helps to cancel the $U(1)_X$ D-term.

In section 2, we analyze in detail this model with global supersymmetry. It is shown, using an effective lagrangian approach, that global supersymmetry is dynamically broken.

In section 3, we compute the resulting soft breaking terms in the observable sector and discuss their phenomenological consequences. We end with some comments.

2 Gauge group $G = SU(N_c) \times U(1)_X$ with $N_f \leq N_c$ flavors

The model that we consider is an extension of SUSY-QCD based on the gauge group $SU(N_c)$ with $N_f \leq N_c$ flavors of "quarks" Q^i of $U(1)_X$ charge q in the fundamental of $SU(N_c)$ and "antiquarks" $\tilde{Q}_{\bar{\imath}}$ of charge $\tilde{q}$ in the antifundamental of $SU(N_c)$.

Since we want to avoid $SU(N_c)$ breaking in the $U(1)_X$ flat direction (1.5), we require that the charges q and $\tilde{q}$ are positive (this is in fact not restrictive: see the comment in the footnote below). We then need at least one field of negative charge in order to cancel the D-term (1.5). For simplicity we will introduce a single field ϕ of $U(1)_X$ charge normalized to -1.

The classical lagrangian compatible with the symmetries is $\mathcal{L} = \mathcal{L}_{kin} + \mathcal{L}_{couplings}$, where we assume a flat Kähler potential for the matter fields:

$$\mathcal{L}_{kin} = \int d^4\theta \left[Q^+ e^{2qV_X + V_N} Q + \tilde{Q} e^{2\tilde{q}V_X - V_N} \tilde{Q}^+ + \phi^+ e^{-2V_X} \phi \right] + \mathcal{L}_{S,V} \quad (2.8)$$

and

$$\mathcal{L}_{couplings} = \int d^2\theta (\frac{\phi}{M_P})^{q+\tilde{q}} m_i^{\bar{\imath}} Q^i \tilde{Q}_{\bar{\imath}} + h.c. \ . \quad (2.9)$$

As stressed in the introduction, the model can be studied *per se* or be used as an illustrative example of a hidden sector where supersymmetry is dynamically broken in presence of an anomalous $U(1)$. In the former case, M_P is the scale of the underlying non-anomalous theory (say the mass of some heavy fermions we have integrated upon), in the latter case it is the Planck scale.

The mixed anomaly $U(1)_X[SU(N_c)]^2$ which will fix, through (1.3), all the mixed anomalies in the model is given by

$$C_N = \frac{1}{4\pi^2} N_f (q + \tilde{q}) = k_N \delta_{GS}. \quad (2.10)$$

We thus require $q + \tilde{q} > 0$, which in turn justifies the presence of the superpotential term (2.9).[3]

[3] Alternatively, since from (2.10) $q + \tilde{q}$ and δ_{GS} have the same sign, we would still avoid

One may note, using (1.2), that the following combination

$$f = k_N S - \frac{N_f}{8\pi^2}(q+\tilde{q}) \ln \frac{\phi}{M_P} \tag{2.11}$$

is invariant under $U(1)_X$. Such a gauge kinetic function would be obtained by integrating over the hidden matter degrees of freedom, assuming unbroken supersymmetry. It could then be used to determine the gaugino masses. We will see however that supersymmetry is broken, which makes matters less straightforward.

The two scales present in the problem are:

- the scale at which the anomalous $U(1)_X$ symmetry is broken which is set by

$$\xi = \frac{1}{2} k_X^{1/2} g_X \delta_{GS}^{1/2} M_P. \tag{2.12}$$

- the scale at which the gauge group $SU(N_c)$ enters in a strong coupling regime:

$$\Lambda = M_P e^{-8\pi^2 k_N S/(3N_c - N_f)}, \tag{2.13}$$

where we have used (1.7) with $b_0 = (3N_c - N_f)/(16\pi^2)$. Notice that, by using the transformation (1.2), we find that the dynamical scale Λ has a charge $q_\Lambda = N_f(q+\tilde{q})/(3N_c - N_f)$.

From now on, we will suppose that $\Lambda << \xi$. We could write the effective theory below the scale ξ and study within this theory the strongly coupled $SU(N_c)$ theory. It is however simpler to keep the complete theory down to the scale Λ since most of the nontrivial effects we will obtain result from an interplay between the scales Λ and ξ.

Below the scale Λ the appropriate degrees of freedom for $N_f < N_c$ are the field ϕ and the mesons $M^i_{\bar{\imath}} = Q^i \tilde{Q}_{\bar{\imath}}$. The effective superpotential is fixed uniquely by the global symmetries [10, 11] as follows

$$W = (N_c - N_f) \frac{\Lambda^{\frac{3N_c - N_f}{N_c - N_f}}}{(detM)^{\frac{1}{N_c - N_f}}} + (\frac{\phi}{M_P})^{q+\tilde{q}} m^{\bar{\imath}}_i M^i_{\bar{\imath}} \tag{2.14}$$

$SU(N_c)$ breaking in the $U(1)_X$ flat direction with $q + \tilde{q} < 0$. The field ϕ which cancels the D-term would then be chosen with charge $+1$.

and is seen to be automatically $U(1)_X$ invariant. Similarly, the gaugino condensation scale

$$< \lambda\lambda > = \left(\Lambda^{3N_c-N_f}/detM\right)^{\frac{1}{N_c-N_f}} \tag{2.15}$$

is also $U(1)_X$ gauge invariant, as it should be.

The gauge contributions to the scalar potential can be computed along the $SU(N_c)$ classical flat directions. The result is

$$V_D = \frac{g_X^2}{2}\left[(q+\tilde{q})Tr(M^+M)^{1/2} - \phi^+\phi + \xi^2\right]^2 \ . \tag{2.16}$$

The auxiliary fields, computed from (2.14) and (2.16) are

$$\bar{F}_{S^+} = -\frac{8\pi^2}{M_P}k_N(S+S^+)^2\frac{\Lambda^{\frac{3N_c-N_f}{N_c-N_f}}}{(detM)^{\frac{1}{N_c-N_f}}} \tag{2.17}$$

and

$$\begin{aligned}
(\bar{F}_{M^+})_i^{\bar{i}} &= 2\left[-(M^{-1})_i^{\bar{j}}\frac{\Lambda^{\frac{3N_c-N_f}{N_c-N_f}}}{(detM)^{\frac{1}{N_c-N_f}}} + (\frac{\phi}{M_P})^{q+\tilde{q}}m_i^{\bar{j}}\right][(M^+M)^{1/2}]_{\bar{j}}^{\bar{i}}, \\
\bar{F}_{\phi^+} &= \frac{q+\tilde{q}}{M_P}(\frac{\phi}{M_P})^{q+\tilde{q}-1}Tr(mM)\ , \\
D_X &= g_X^2\left[(q+\tilde{q})Tr(M^+M)^{1/2} - \phi^+\phi + \xi^2\right] \ .
\end{aligned} \tag{2.18}$$

In the limit $S \to \infty$, the scale Λ vanishes and we can choose $M_{\bar{i}}^i = 0$ and $\phi = \xi$ to cancel all these auxiliary fields. This is the usual global supersymmetry minimum at infinite values of S which leads to the dilaton stabilization problem. We will assume that the dilaton is stabilized at some finite value S_0, possibly through some extra S-dependent term in the superpotential (we will therefore refrain from using (2.17)), and that $F_S(S_0) = 0$. Indeed we are going to show that, even in this unfavorable case (supersymmetry conserving groundstate for S), the other fields present in the theory yield supersymmetry breaking because of the anomalous behavior of $U(1)_X$. From now on, we will therefore restrict our attention to the auxiliary fields (2.18) associated with $M_{\bar{i}}^i$, ϕ and the $U(1)_X$ gauge degree of freedom.

It is readily seen that the system of equations $F_M = F_\phi = D_X = 0$ has no solution as long as $\xi \neq 0$ that is $q + \tilde{q} \neq 0$. Therefore supersymmetry is dynamically broken. As usual, the origin of supersymmetry breaking is chiral: the non-abelian gauge group content is vector-like but Q^i and $\tilde{Q}_{\bar{\imath}}$ do not transform in a vectorlike fashion under $U(1)_X$: $q \neq -\tilde{q}$; this is precisely what drives the anomaly, which is therefore the source of chirality in the model.

We will now minimize the scalar potential in terms of $M^i_{\bar{\imath}}$, ϕ at fixed $S = S_0$ value such that $F_S(S_0) = 0$:

$$V = \frac{1}{(S+S^+)^2}\bar{F}_{S^+}F_S + \bar{F}_{\phi^+}F_\phi + \frac{1}{2}(\bar{F}_{M^+})^{\bar{\imath}}_i[(M^+M)^{-1/2}]^{\bar{j}}_{\bar{\imath}}(F_M)^i_{\bar{j}} + V_D. \tag{2.19}$$

In order to be able to give analytic solutions, we make a few simplifying asumptions. First, we linearize the minimization procedure by looking for a minimum in the vicinity of:

a) $\phi_0 = \xi$, the field value which minimizes V_D in the absence of condensates;

b) $(M_0)^i_{\bar{\imath}}$, the solution of $(\bar{F}_{M^+})^{\bar{\imath}}_i = 0$:

$$(M_0)^i_{\bar{\imath}} = (m^{-1})^i_{\bar{\imath}}(detm)^{1/N_c}\Lambda^{\frac{3N_c-N_f}{N_c}}\left(\frac{\xi}{M_P}\right)^{\frac{N_f-N_c}{N_c}(q+\tilde{q})} \quad . \tag{2.20}$$

One can make a field transformation in order to have a diagonal matrix $m^{\bar{\imath}}_i$, in which case $(M_0)^i_{\bar{\imath}}$ is also diagonal. Since we are only interested in orders of magnitude, we will make the assumption that $m^{\bar{\imath}}_i = m\delta^{\bar{\imath}}_i$ and search for solutions $M^i_{\bar{\imath}} = M\delta^i_{\bar{\imath}}$ of the equations of motion.

The minimum is obtained by making around the field configuration $(M_0)^i_{\bar{\imath}} \equiv M_0\delta^i_{\bar{\imath}}$ an expansion in the parameter

$$\epsilon \equiv \frac{M_0}{\xi^2} = \left(\frac{\Lambda}{\xi}\right)^{\frac{3N_c-N_f}{N_c}}\left[\frac{m}{M_P}\left(\frac{\xi}{M_P}\right)^{q+\tilde{q}-1}\right]^{\frac{N_f-N_c}{N_c}} \quad . \tag{2.21}$$

One obtains

$$<\phi^+\phi> \;=\; \xi^2\left[1+\epsilon N_f(q+\tilde{q}) + \epsilon^2 N_f^2(q+\tilde{q})^2\left(-\frac{(N_c-N_f)(2N_c-2N_f)}{2N_c^2}(q+\tilde{q})\right.\right.$$

$$+ \quad \frac{1}{g_X^2}\hat{m}^2\left[1-\frac{N_f}{N_c}(q+\tilde{q})\right]\Bigg) + O(\epsilon^3)\Bigg], \tag{2.22}$$

$$< M > \;=\; M_0\left[1-\frac{\epsilon}{2}\frac{N_f(N_c-N_f)(2N_c-N_f)}{N_c^2}(q+\tilde{q})^2 + O(\epsilon^2)\right], \tag{2.23}$$

where we have introduced the scale

$$\hat{m} = m\left(\frac{\xi}{M_P}\right)^{q+\tilde{q}} \quad . \tag{2.24}$$

The value of the auxiliary terms at this ground state are, to leading order:

$$\begin{aligned} < D_X > &= -\epsilon^2\hat{m}^2 N_f^2 (q+\tilde{q})^2\left[1-\frac{N_f}{N_c}(q+\tilde{q})\right], \\ < F_\phi > &= \epsilon\hat{m}\xi N_f(q+\tilde{q}), \\ < F_M > &= -\epsilon^2\hat{m}\xi^2\frac{N_f(N_c-N_f)}{N_c}(q+\tilde{q})^2. \end{aligned} \tag{2.25}$$

One may note that $< D_X^{1/2} >$, $< F_\phi/\phi >$ and $< F_M/M >$ are all of the same order $\epsilon\tilde{m}$. This will have definite consequences for the soft terms, as we will see in the next section. Also, using (2.21), one checks that this order of magnitude goes to zero as $\Lambda \to 0$, as well as $\xi \to \infty$ (at fixed value of ξ/M_P). This shows once again the mixed role that the two scales Λ and ξ play as far as supersymmetry breaking is concerned in this model.

A similar analysis can be performed when $N_f = N_c \equiv N$. In this case, new degrees of freedom must be introduced in the low energy effective lagrangian [12]

$$B = \epsilon_{i_1\cdots i_N} Q^{i_1}\cdots Q^{i_N}, \;\; \tilde{B} = \epsilon^{\bar{\imath}_1\cdots\bar{\imath}_N}\tilde{Q}_{\bar{\imath}_1}\cdots Q_{\bar{\imath}_N}. \tag{2.26}$$

The effective superpotential compatible with all the symmetries reads

$$W = U\ln\frac{detM - B\tilde{B}}{\Lambda^{2N}} + \left(\frac{\phi}{M_P}\right)^{q+\tilde{q}} m_i^{\bar{\imath}} M_{\bar{\imath}}^i, \tag{2.27}$$

where U is a Lagrange multiplier, physically interpreted as the gauge composite superfield $U = \mathrm{Tr}W_N^\alpha W_{N\alpha}$. As in the $N_f < N_c$ case, the system of equations $F_U = F_M = F_B = F_{\tilde{B}} = D_X = 0$ has no solution and global supersymmetry is broken.

The properties of the model do not depend either on the assumption of a single ϕ field. Similar conclusions can be reached when one considers for example a vector-like pair of such fields.

3 Soft terms in the observable sector

We now use the results of the preceding section to determine the order of magnitude of the soft terms in the observable low energy sector of quarks, leptons and gauginos.

The magnitude of the soft terms in the observable sector is fixed by the auxiliary fields F_ϕ, F_M and D_X. At the tree level of the Lagrangian of our model, we find soft scalar masses $\tilde{m}_i^2$ and trilinear soft terms A_{ijk} given by the expressions[4]:

$$\tilde{m}_i^2 = X_i < D_X >, \quad A_{ijk} = (X_i + X_j + X_k)\frac{F_\phi}{\phi}, \tag{3.28}$$

where X_i is the $U(1)_X$ charge of the corresponding field Φ^i. Gaugino masses in the hidden sector are also induced: $m_\lambda \sim N_f < F_M/M >$. The gaugino masses in the observable sector are absent at tree level and are induced by standard gauge loops. Notice that, because supersymmetry is broken, gaugino masses are not simply given by the gauge invariant kinetic function (2.11).

In the preceding section we obtained, in the limit $\Lambda \ll \xi$, the following relation among the auxiliary fields $< F_M/M > \sim < F_\phi/\phi > \sim < D_X^{1/2} >$. Consequently, all the soft breaking terms induced at tree level are of the same order. By using (2.21) and (2.25), we obtain

$$\tilde{m} \sim N_f(q+\tilde{q})\frac{\Lambda^3}{\xi^2}\left[\frac{m}{\Lambda}\left(\frac{\xi}{M_P}\right)^{q+\tilde{q}}\right]^{N_f/N_c} = N_f(q+\tilde{q})\frac{<\lambda\lambda>}{\xi^2}, \tag{3.29}$$

where $\tilde{m}$ generically denotes a soft-breaking term (3.28) and we have used (2.15) in order to derive the last relation. This relation is indeed central

[4] We assume the presence in the superpotential of terms of the form $(\phi/M_P)^{X_i+X_j+X_k}\Phi_i\Phi_j\Phi_k$, as allowed by the $U(1)_X$ symmetry.

to the kind of models described here and stresses the connected role of the relevant scales: ξ as the scale of messenger interaction and the gaugino condensate as the seed of supersymmetry breaking (although, as stressed earlier, the chiral nature of the $U(1)_X$ plays an important role: $q \neq -\tilde{q}$).

We now restrict our attention to hidden sector models where the messengers of supersymmetry breaking are the anomalous $U(1)_X$ gauge *and* gravitational interactions. The scale M_P is therefore the Planck scale. Eq. (3.29) should be compared with the gravitationally-induced soft terms of order

$$\tilde{m}|_{grav.} \sim \frac{<\lambda\lambda>}{M_P^2}, \tag{3.30}$$

which must be included if the supergravity interactions are switched on. Since

$$\frac{1}{N_f(q+\tilde{q})}\left(\frac{\xi}{M_P}\right)^2 \tag{3.31}$$

is a small number in the context of superstring models, the supergravity soft terms can be neglected and the soft terms computed above can be viewed as phenomenological predictions of these models. It is worth noting that, using (2.10) and (2.12), the factor $N_f(q+\tilde{q})$, which may be large and is model dependent, drops out out of the ratio (3.31). One is left with a pure number which depends on the gauge coupling –as a genuine one-loop effect– and the Kac-Moody levels.

If $\xi \sim M_P$, phenomenologically interesting soft terms $\tilde{m} \leq 1$ TeV call for $\Lambda \ll \Lambda_0$, where Λ_0 is a typical intermediate condensation scale: $\Lambda_0 \sim 10^{13}$ TeV.[5] For $\xi \ll M_P$ the required values of Λ depend on the parameters of the hidden sector N_f, N_c, $q+\tilde{q}$. A generic prediction in this case is a rather light gravitino:

$$m_{3/2} \sim \frac{<W>}{M_P^2} \sim \frac{N_c}{N_f(q+\tilde{q})}\left(\frac{\xi}{M_P}\right)^2 \tilde{m}. \tag{3.32}$$

It is useful to notice that generally $N_f(q+\tilde{q}) > N_c$. Then, by using (2.25), we find $D_X > 0$, which means that scalar particles with positive $U(1)_X$ charges

[5] We assume here that $m \sim M_P$.

will acquire positive squared masses. This is a very welcome feature of the model since, by using the mixed anomaly conditions (1.3), (1.4), we know that the fields in the observable sector must have predominantly positive charges.

This is also consistent with the interpretation of the $U(1)_X$ as an horizontal symmetry which may explain the low energy mass spectrum [5, 13, 6, 14, 15]. In this context, positive charges are required in the model presented above to account for the observed fermion masses and mixings. The anomalous symmetry also plays an important role in constraining the soft terms [16, 17, 18].

A general feature of the models presented is the large scale of gauge symmetry breaking. Despite the fact that the messengers of supersymmetry breaking are mostly gauge interactions, the resulting soft terms have much lower values, as a result of a conspiracy between the two scales Λ and ξ. This is to be contrasted with the standard gauge-mediated scenarios where the scale of supersymmetry breaking is much lower [19]. In particular, as we have seen, in the limit where the two scales are very far away $\Lambda \to 0$ or $\xi \to \infty$, supersymmetry is restored.

There are several aspects which we did not address in this short paper and which we reserve for further studies: one is the question of dilaton stabilization since the scalar potential has a complicate S dependence in its F-term through the scale Λ and its D-term through the coupling g_X. Another is the generalization to other gauge structures with a chiral content and explicit superstring realizations.

Acknowledgments: We wish to thank the Aspen Center for Physics where this work was completed for its hospitality and the participants of the workshop "Flavor and Gauge Hierarchies", in particular Pierre Ramond, for the lively discussions and the interesting questions they raised.

References

[1] H.P. Nilles, *Phys. Lett.* **B115** (1982) 193; S. Ferrara, L. Girardello and H.P. Nilles, *Phys. Lett.* **B125** (1983) 457.

[2] J.P. Derendinger, L.E. Ibáñez and H.P. Nilles, *Phys. Lett.* **B155** (1985) 65; M Dine, R. Rohm, N. Seiberg and E. Witten, *Phys. Lett.* **B156** (1985) 55.

[3] N. V. Krasnikov, *Phys. Lett.* **B193** (1987) 37.

[4] M. Dine, N. Seiberg and E. Witten, *Nucl. Phys.* **B289** (1987) 589; J. Atick, L. Dixon and A. Sen, *Nucl. Phys.* **B292** (1987) 109; M. Dine, I. Ichinose and N. Seiberg, *Nucl. Phys.* **B293** (1988) 253.

[5] L. Ibáñez and G.G. Ross, *Phys. Lett.* **B332** (1994) 100.

[6] P. Binétruy and P. Ramond, *Phys. Lett.* **B350** (1995) 49; P. Binétruy, S. Lavignac and P. Ramond, LPTHE-Orsay 95/54, UFIFT-HEP-96-1, hep-ph/9601243 (to be published in Nuclear Physics).

[7] J.A. Casas and C. Munoz, *Phys. Lett.* **B216** (1989) 37; J.C. Casas, J. M. Moreno, C. Munoz and M. Quiros, *Nucl. Phys.* **B328** (1989) 272; P. Binétruy and G. Dvali, preprint CERN-TH/96-149, LPTHE-ORSAY 96/40.

[8] T. Banks and M. Dine, preprint RU-95-51, SCIPP 95/41.

[9] M. Green and J. Schwarz, *Phys. Lett.* **B149** (1984) 117.

[10] T.R. Taylor, G. Veneziano and S. Yankielowicz, *Nucl. Phys.* **B218** (1983) 493.

[11] I. Affleck, M. Dine and N. Seiberg, *Nucl. Phys.* **B256** (1985) 557.

[12] N. Seiberg, *Phys. Rev.* **D49** (1994) 6857.

[13] V. Jain and R. Schrock, *Phys. Lett.* **B352** (1995) 83.

[14] E. Dudas, S. Pokorski and C.A. Savoy, *Phys. Lett.* **B356** (1995) 45.

[15] Y. Nir, *Phys. Lett.* **B354** (1995) 107.

[16] M. Leurer, Y. Nir and N. Seiberg, *Nucl. Phys.* **B398** (1993) 319 and **B309** (1993) 337.

[17] E. Dudas, S. Pokorski and C.A. Savoy, *Phys. Lett.* **B369** (1996) 255; E. Dudas, C. Grojean, S. Pokorski and C.A. Savoy, Saclay T96/065, hep-ph/9606383.

[18] Y. Kawamura and T. Kobayashi, DPSU-96-2, hep-ph/9601365.

[19] M. Dine and A. Nelson, *Phys. Rev.* **D48** (1993) 1277; M. Dine, A. Nelson and Y. Shirman, *Phys. Rev.* **D51** (1995) 1362; M. Dine, A. Nelson, Y. Nir and Y. Shirman, *Phys. Rev.* **D53** (1996) 2658.